Eur. Phys. J. C (2014) 74:2853
DOI 10.1140/epjc/s10052-014-2853-z

THE EUROPEAN PHYSICAL JOURNAL C

Regular Article - Theoretical Physics

On neutrinoless double beta decay in the minimal left–right symmetric model

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Received: 31 January 2014 / Accepted: 9 April 2014 / Published online: 6 May 2014

Abstract We analyze the general phenomenology of neutrinoless double beta decay in the minimal left–right symmetric model. We study under which conditions a New Physics dominated neutrinoless double beta decay signal can be expected in the future experiments. We show that the correlation among the different contributions to the process, which arises from the neutrino mass generation mechanism, can play a crucial role. We have found that, if no fine tuned cancelation is involved in the light–active neutrino contribution, a New Physics signal can be expected mainly from the W_R–W_R channel. An interesting exception is the W_L–W_R channel which can give a dominant contribution to the process if the right-handed neutrino spectrum is hierarchical with $M_1 \lesssim$ MeV and $M_2, M_3 \gtrsim$ GeV. We also discuss if a New Physics signal in neutrinoless double beta decay experiments is compatible with the existence of a successful Dark Matter candidate in the left–right symmetric models. It turns out that, although it is not a generic feature of the theory, it is still possible to accommodate such a signal with a KeV sterile neutrino as dark matter.

1 Introduction

The recent LHC results [1,2] seem to indicate that the Higgs mechanism, with the Higgs mass around 125 GeV, is the responsible for the mass generation of the Standard Model (SM) particles. However, the origin of light neutrino masses, for the existence of which we have compelling evidence from neutrino oscillation experiments, still remains unknown. It is true that the light neutrino masses could also be generated through the Higgs mechanism in a minimally extended SM which includes sterile (right-handed) neutrino fields as $SU(2)_L$ singlets and in which the total lepton number is conserved. However, their smallness in comparison with the charged lepton and quark masses calls for a different explanation. In this context, extensions of the SM required to explain the origin of neutrino masses, and compatible with the latest LHC data, arise as quite suggestive models of New Physics (NP). Among those we find the celebrated seesaw models [3–6], which can give, in addition, the key to explain the matter-antimatter asymmetry of the universe through leptogenesis [7].

Most of those models predict that neutrinos are Majorana particles, something which can be tested in lepton number violating processes such as the neutrinoless double beta ($0\nu\beta\beta$) decay. The $0\nu\beta\beta$ decay experiments are the most promising ones in this context but they suffer a serious drawback: the NP contribution to the process is usually short range and thus typically very suppressed compared to that of the light neutrinos. Thanks to the future $0\nu\beta\beta$ experiments [8–14], in combination with the complementary information coming from neutrino oscillation experiments and cosmology, we might be able to discover the Majorana nature of neutrinos, but not easily which is the mechanism responsible for the neutrino mass generation [15,16]. In this context, the correlations between the standard light neutrino and NP contribution to the $0\nu\beta\beta$ decay are crucial, as shown in the case of the type-I [3–6], type-II [17–21] and type-III [22] seesaw models in Refs. [23,24]. The generation of light neutrino masses in a particular model usually induces important correlations between the different contributions to the $0\nu\beta\beta$ decay, which should always be considered in a model dependent analysis, helping to understand which type of NP it is feasible to test in the experiments.

In this work we will focus on the $0\nu\beta\beta$ decay phenomenology of the minimal left–right symmetric model (MLRSM) [6,21,25–27]. The left–right symmetric models have been widely studied in the literature since, among other features,

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they provide a natural explanation for the smallness of the neutrino masses (some recent analysis in the context of the $0\nu\beta\beta$ decay can be found in Refs. [28–35]). In our analysis we will assume that no accidental cancelation occurs in the light neutrino mediated W_L–W_L channel, which involves the exchange of two W_L. We will distinguish three regions of the parameter space depending on the mass of the right-handed (RH) neutrinos. First, we will show that if the right-handed (RH) neutrinos are heavier than the $0\nu\beta\beta$ decay scale ($\sim$100 MeV), the $0\nu\beta\beta$ decay rate is dominated by light neutrino exchange channels with the exception of the channel in which two W_R are exchanged (W_R–W_R channel) mediated by heavy neutrinos[1]. One of the light neutrino mediated channels involves the exchange of one W_L and one W_R (W_L–W_R channel); however, it turns out that a NP dominant contribution can come mainly from the W_R–W_R channel. Secondly, we will study the region of the parameter space where the RH neutrinos are lighter than the $0\nu\beta\beta$ scale. We have found that in this case the W_L–W_R contribution cancels out while a NP signal can still be expected from the W_R–W_R channel. In this region, the RH neutrinos can give a relevant contribution through the W_L–W_L channel, as opposed to the type-I seesaw case where the total W_L–W_L contribution is very suppressed. Finally, we will investigate a mixed scenario with RH neutrinos in both regions below and above the $0\nu\beta\beta$ decay scale. We have found that this is the only scenario in which the W_L–W_R channel turns out to be relevant and can be responsible of a future signal (if no cancelation in the W_L–W_L channel is invoked). In all the cases we will show for which part of the parameter space a NP signal in future $0\nu\beta\beta$ decay experiments can be expected. Moreover, we will also analyze if such a signal can be compatible with the existence of a successful Dark Matter (DM) candidate in the left–right symmetric model, study the complementary bounds coming from charged lepton flavor violation (LFV) experiments and the impact of the one-loop corrections to the light neutrino masses.

This work is organized as follows. In Sect. 2 we briefly describe the MLRSM, focusing on the relations among the parameters of the model induced by the neutrino mass generation. In Sect. 3 we analyze the neutrinoless double beta decay phenomenology in the MLRSM, studying in particular for which part of the parameter space a $0\nu\beta\beta$ decay signal coming mainly from NP contributions can be possible. Section 4 is devoted to the analysis of complementary constraints coming mainly from charged LFV experiments and the stability of the light neutrino masses under one-loop corrections. In Sect. 5 we study if a successful DM candidate and a NP signal in the future $0\nu\beta\beta$ decay experiments can be compatible in the MLRSM. Finally, we conclude in Sect. 6.

[1] If other contributions coming from different models are not involved.

2 Minimal left–right symmetric model and neutrino masses

The Lagrangian of the MLRSM respects an enlarged gauge symmetry $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ plus a discrete left–right symmetry which leads to equal $SU(2)_L$ and $SU(2)_R$ gauge couplings ($g_L = g_R = g$). We are not going into the details of the model since it has been widely studied in the literature (for a recent complete analysis regarding the associated lepton number violating effects, see for instance Refs. [29,33]), but only recall the most relevant features for our analysis. The scalar sector is also augmented by the addition of two scalar triplets (Δ_L and Δ_R) and a bi-doublet scalar under $SU(2)_L \otimes SU(2)_R$, which spontaneously break the electroweak symmetry when they develop vacuum expectation values (VEVs).

In this section we will derive the relations which will be used in the phenomenological analysis of the $0\nu\beta\beta$ decay. Since they come from the neutrino mass generation, let us recall how the complete neutrino mass matrix looks like after the electroweak symmetry breaking:

$$M_\nu = \begin{pmatrix} M_L & m_D^T \\ m_D & M_R \end{pmatrix} = U \,\mathrm{Diag}(m, M) U^T, \tag{1}$$

where m_i are the light neutrino masses and M_i the heavy ones. Notice that in this model the Majorana mass term for the heavy neutrinos is generated dynamically when Δ_R takes a VEV ($M_R = Y_{\Delta_R} v_R$), while the Majorana mass term M_L for the left-handed (LH) neutrinos is generated analogously through the Δ_L VEV ($M_L = Y_{\Delta_L} v_L$). The neutrino mass matrix is diagonalized as shown above by a 6×6 unitary matrix U, through the following rotation between the neutrino flavor and mass eigenstates denoted by $\alpha, \beta = e, \mu, \tau$ and $i, k = 1, 2, 3$, respectively:

$$\begin{pmatrix} \nu_{\alpha L} \\ N^c_{\beta R} \end{pmatrix} = U \begin{pmatrix} \nu_i \\ N_k \end{pmatrix} = \begin{pmatrix} \tilde{U} & B \\ A & V \end{pmatrix} \begin{pmatrix} \nu_i \\ N_k \end{pmatrix}. \tag{2}$$

The diagonalization of the complete neutrino mass matrix presented in Eq. (1) provides the following useful relations:

$$\tilde{U} m \tilde{U}^T + B M B^T = M_L, \tag{3}$$
$$\tilde{U} m A^T + B M V^T = m_D^T, \tag{4}$$
$$A m A^T + V M V^T = M_R. \tag{5}$$

On the other hand, taking into account that the active LH block of U, $\tilde{U}$, is unitary to a very good approximation (at least up to the percent level [36]), the complete neutrino mixing matrix can be expanded as

$$\begin{aligned} U &= \begin{pmatrix} 1 - \theta\theta^\dagger/2 & \theta \\ -\theta^\dagger & 1 - \theta^\dagger\theta/2 \end{pmatrix} \begin{pmatrix} U_{\text{pmns}} & 0 \\ 0 & V \end{pmatrix} + \mathcal{O}\left(\theta^3\right) \\ &= \begin{pmatrix} U_{\text{pmns}} & \theta V \\ -\theta^\dagger U_{\text{pmns}} & V \end{pmatrix} + \mathcal{O}\left(\theta^2\right), \end{aligned} \tag{6}$$

where θ is a 3×3 matrix which characterizes the small mixing between the active LH and the heavy RH neutrinos, $U_{\rm pmns}$ is the PMNS matrix and V is a 3×3 unitary matrix. From Eqs. (3–5), we have

$$U_{\rm pmns}\, m\, U^T_{\rm pmns} = M_L - \theta M_R \theta^T, \tag{7}$$

$$\theta M_R - M_L \theta^* = m_D^T, \tag{8}$$

$$V M V^T = M_R \left(1 + \mathcal{O}\left(\frac{M_L}{M_R}\theta^2\right)\right). \tag{9}$$

The discrete (charge conjugation) LR symmetry gives the following relation between the Yukawa couplings of the triplets: $Y_{\Delta_R} = Y_{\Delta_L} \equiv Y_\Delta$.[2] This means that

$$(M_L)_{\alpha\beta}/(M_R)_{\alpha\beta} = v_L/v_R < 10^{-3}, \tag{10}$$

where we have employed the present bounds on v_L and v_R, namely $v_L \lesssim 7$ GeV [37] and $v_R \gtrsim 10$ TeV ($M_{W_R} \approx g v_R/\sqrt{2} \gtrsim$ TeV [38–40]). Therefore, the $\mathcal{O}(M_L\theta^*)$ and $\mathcal{O}\left(\frac{M_L}{M_R}\theta^2\right)$ can be safely neglected in Eq. (8) and Eq. (9), respectively, and

$$\theta \simeq m_D^T M_R^{-1}. \tag{11}$$

Of course, θ plays a fundamental role at the phenomenological level since it basically describes the mixing between the active LH neutrinos and the RH ones. It would be very interesting thus to find a useful parameterization of θ as a function of the light neutrino parameters, light neutrino masses and the angles/phases of the PMNS matrix, and the rest of the independent parameters of the model associated with the RH neutrino sector. In principle, an analogous parameterization to the Casas–Ibarra one [41] would be a good candidate [42]. However, the presence of M_L in Eq. (7) and the fact that the matrix V is in this case physical, contrary to the type-I seesaw model, makes that parameterization less transparent and more involved than expected. On the other hand, the discrete (charge conjugation) LR symmetry leads to the following constraint:

$$m_D = m_D^T, \tag{12}$$

and thus Eq. (11) becomes

$$\theta M_R = M_R \theta^T = m_D. \tag{13}$$

Plugging this relation and Eq. (9) into Eq. (7), we obtain,

$$U_{\rm pmns}\, m\, U^T_{\rm pmns} = M_L - \theta^2 V M V^T, \tag{14}$$

and finally with $Y_{\Delta_R} = Y_{\Delta_L} \equiv Y_\Delta$ and hence $M_L = \frac{v_L}{v_R} M_R$, we have

$$\theta = \left[\frac{v_L}{v_R} I - U_{\rm pmns}\, m\, U^T_{\rm pmns} V^* M^{-1} V^\dagger\right]^{1/2}. \tag{15}$$

[2] Another option is to consider instead a discrete parity symmetry leading to a similar relation: $Y_{\Delta_R} = Y^*_{\Delta_L} \equiv Y_\Delta$.

Therefore, θ is completely determined as a function of the light and heavy neutrino masses, m and M, the PMNS matrix, $U_{\rm pmns}$, v_L/v_R, and the unitary matrix V [31]. Notice that if this expression is used to obtain θ with the PMNS mixing angles and the solar and atmospheric mass-squared differences as input parameters, we ensure that the model is consistent with the light neutrino mass and mixing pattern measured in neutrino oscillation experiments.

3 Neutrinoless double beta decay

In our study of the $0\nu\beta\beta$ decay in the MLRSM, we will pay special attention to the correlation among all the contributions to the process and, in particular, the connection with the light neutrino masses. We shall see that the correlation between the different contributions and the experimental bounds on the parameters will allow us to safely neglect some of the NP contributions.

As we have already mentioned, we will not analyze the scenario in which a cancelation occurs within the standard light neutrino contribution, which would naively leave the NP channels as the leading contributions [43,44]. Of course, this cancelation can be due to the presence of an extra symmetry added to the model, such as the lepton number which is approximately conserved in the so called inverse or direct seesaw models [45–49]. The problem in this scenario is that, in order for the NP contributions to be measurable, a significant violation of lepton number should be introduced through the NP sector which may not have an impact on the light neutrino masses at tree level but arises naturally at one-loop level, as shown in Ref. [23] in the context of the seesaw models. This makes it very difficult to have a significant contribution from NP channels since the one-loop correction to the light neutrino masses tends to dominate in the $0\nu\beta\beta$ decay rate.

We will distinguish three different regions according to the associated $0\nu\beta\beta$ decay phenomenology: (i) when the RH neutrinos are much heavier than the $0\nu\beta\beta$ decay scale ($\langle p\rangle \approx 100$ MeV), which means heavier than approximately 1 GeV; (ii) when the RH neutrinos are much lighter than the $0\nu\beta\beta$ decay scale (below 1 MeV); (iii) when the RH neutrinos are in both regions, (i) and (ii).

In the analysis below we have reasonably estimated the NMEs corresponding to some of the channels under study. This is accurate enough for our purposes but, although the associated NMEs errors are still large, in order to be more precise, the full calculation of all the NMEs should be considered.

3.1 Heavy regime

The various contributions to the $0\nu\beta\beta$ transition rate in this model are described by the Feynman diagrams shown in

Fig. 1 Feynman diagrams contributing to the $0\nu\beta\beta$ transition rate in the MLRSM

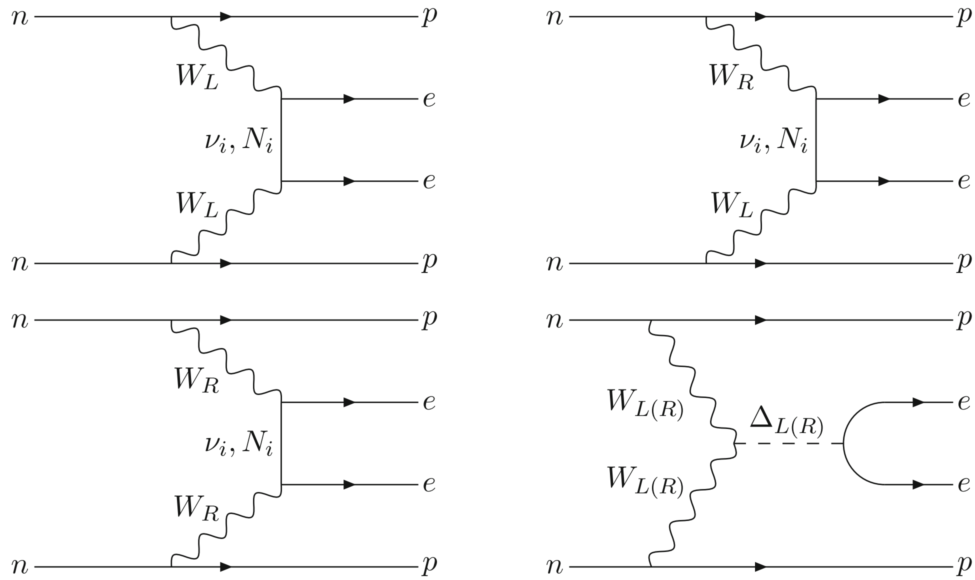

Fig. 1. We will start describing them one by one in order to show that the contributions of the diagrams in which a heavy fermion (or scalar) is exchanged are subdominant with respect to those of the light neutrino exchange, the only exception being the W_R–W_R channel.

- **$\mathbf{W_L - W_L}$ channel**. The amplitude corresponding to the top left diagram of Fig. 1 is given by

$$A_{LL} \propto (1 + \mathcal{O}(\xi)) \left[\sum_{i=1}^{3} m_i \tilde{U}_{ei}^2 \mathcal{M}^{0\nu\beta\beta}(0) + \sum_{i=1}^{3} M_i B_{ei}^2 \mathcal{M}^{0\nu\beta\beta}(M_i) \right],$$

where ξ is the W_L–W_R mixing angle. The present experimental bound is given by $\xi \lesssim 10^{-2}$ [38], however, in the minimal left–right symmetric model there is a stronger theoretical upper bound given by $M_{W_L}^2/M_{W_R}^2 < 10^{-3}$ [40].[3] In the case of the W_L–W_L contribution, ξ can be safely neglected. $\mathcal{M}^{0\nu\beta\beta}$ are the associated nuclear matrix elements (NMEs) following the notation of Ref. [24], where the NMEs were computed as a function of the mass of the neutrino mediating the process for different nuclei. Notice that in this notation the NMEs include the dependence on the propagator. The NMEs corresponding to the light neutrino exchange are independent of the neutrino masses, and then with Eq. (3), the above amplitude can be rewritten as

$$A_{LL} \propto (M_L)_{ee} \mathcal{M}^{0\nu\beta\beta}(0) + \sum_{i=1}^{3} M_i B_{ei}^2 \left(\mathcal{M}^{0\nu\beta\beta}(M_i) - \mathcal{M}^{0\nu\beta\beta}(0) \right). \qquad (16)$$

Taking into account that $\mathcal{M}^{0\nu\beta\beta}(M_i)/\mathcal{M}^{0\nu\beta\beta}(0) \ll 1$ [24], the contribution due to the heavy neutrino exchange can be safely neglected. Using again Eq. (3), one obtains,

$$A_{LL} \propto \sum_{i=1}^{3} m_i \tilde{U}_{ei}^2 \mathcal{M}^{0\nu\beta\beta}(0) = \sum_{i=1}^{3} \left(U_{\rm pmns} m U_{\rm pmns}^T \right)_{ee} \mathcal{M}^{0\nu\beta\beta}(0), \qquad (17)$$

which is the standard light neutrino contribution.

- **$\mathbf{W_R - W_R}$ channel**. The amplitude of the bottom left diagram of Fig. 1, in which two W_R are involved, is given by [50]

$$A_{RR} \propto \left(\frac{M_{W_L}^2}{M_{W_R}^2} + \xi \right)^2 \left[\sum_{i=1}^{3} m_i A_{ei}^{*2} \mathcal{M}^{0\nu\beta\beta}(0) + \sum_{i=1}^{3} M_i V_{ei}^{*2} \mathcal{M}^{0\nu\beta\beta}(M_i) \right].$$

Using Eq. (6) in the above equation we obtain

$$A_{RR} \propto \left(\frac{M_{W_L}^2}{M_{W_R}^2} + \xi \right)^2 \left[\sum_{i=1}^{3} M_i V_{ei}^{*2} \mathcal{M}^{0\nu\beta\beta}(M_i) - \left(\theta^T U_{\rm pmns}^* m U_{\rm pmns}^\dagger \theta \right)_{ee} \mathcal{M}^{0\nu\beta\beta}(0) \right]. \qquad (18)$$

[3] Notice that ξ can only saturate this bound if and only if the two VEVs of the Higgs doublets are of the same order.

Clearly, the second term can be neglected in comparison with the standard contribution due to the double suppression coming from $\left(\frac{M_{W_L}^2}{M_{W_R}^2}+\xi\right)^2$ and the active–heavy mixing, at least $|\theta_{\alpha i}|^2 \lesssim 10^{-2}$ [36,51,52]. The first term, however, cannot be neglected, i.e.,

$$A_{RR} \propto \sum_{i=1}^{3} M_i V_{ei}^{*2} \left(\frac{M_{W_L}^2}{M_{W_R}^2}+\xi\right)^2 \mathcal{M}^{0\nu\beta\beta}(M_i). \tag{19}$$

- $\mathbf{W_L - W_R}$ **channel**. For the diagram in the top right of Fig. 1, in which W_L and W_R are exchanged, the amplitude is given by

$$A_{LR} \propto \left(\xi + \eta\frac{M_{W_L}^2}{M_{W_R}^2}\right)\langle p\rangle \left[\sum_{i=1}^{3} A_{ei}\tilde{U}_{ei}^{*}\mathcal{M}^{0\nu\beta\beta}(0) + \sum_{i=1}^{3} V_{ei}B_{ei}^{*}\mathcal{M}^{0\nu\beta\beta}(M_i)\right], \tag{20}$$

where $\eta \approx 10^{-2}$ [53–58][4]. Taking into account that U is unitary, we have

$$A_{LR} \propto \left(\xi + \eta\frac{M_{W_L}^2}{M_{W_R}^2}\right)\langle p\rangle \times \sum_{i=1}^{3} V_{ei}B_{ei}^{*}(\mathcal{M}^{0\nu\beta\beta}(M_i) - \mathcal{M}^{0\nu\beta\beta}(0)), \tag{21}$$

and since $\mathcal{M}^{0\nu\beta\beta}(M_i)/\mathcal{M}^{0\nu\beta\beta}(0) \ll 1$, Eq. (21) becomes

$$A_{LR} \propto \left(\xi + \eta\frac{M_{W_L}^2}{M_{W_R}^2}\right)\langle p\rangle \sum_{i=1}^{3} A_{ei}\tilde{U}_{ei}^{*}\mathcal{M}^{0\nu\beta\beta}(0)), \tag{22}$$

which implies that the light neutrino mediated contribution of the W_L–W_R channel is again dominant over the heavy neutrino exchange.

- The amplitude corresponding to the scalar triplet Δ_L exchange (bottom right in Fig. 1 with W_L and Δ_L) is suppressed with the factor

$$\frac{(M_L)_{ee}}{\sum_i \tilde{U}_{ei}m_i}\frac{\langle p^2\rangle}{M_{\Delta_L}^2} = \frac{(M_L)_{ee}}{\left(M_L - m_D M_R^{-1} m_D^T\right)_{ee}}\frac{\langle p^2\rangle}{M_{\Delta_L}^2}, \tag{23}$$

with respect to the standard contribution given in Eq. (17). The suppression factor is at least $\langle p^2\rangle/M_{\Delta_L}^2 \ll 1$ if no fine tuned cancelation between the two terms in the light neutrino contribution is invoked, i.e., the contribution of this channel is negligible. For the corresponding "right-handed" version of the diagram the situation is slightly different and the suppression factor now reads

$$\frac{M_{W_L}^4}{M_{W_R}^4}\frac{(M_R)_{ee}}{\left(M_L - m_D M_R^{-1} m_D^T\right)_{ee}}\frac{\langle p^2\rangle}{M_{\Delta_R}^2}. \tag{24}$$

It seems that for small enough values of $(M_L - m_D M_R^{-1} m_D^T)_{ee} = \sum_i \left[(U_{\text{pmns}})_{ei}\right]^2 m_i$, this contribution could be larger than the standard one. However, it is not very easy to achieve a measurable Δ_R contribution, at the reach of the sensitivity of the next-to-next of $0\nu\beta\beta$ decay experiments ($m_{\beta\beta} \sim 10^{-2}$ eV). Indeed, the corresponding amplitude is given by

$$A_{\Delta_R} \propto (M_R)_{ee}\,\mathcal{M}_{\Delta}^{0\nu\beta\beta}(M_{\Delta_R}) \approx \frac{M_{W_L}^4}{M_{W_R}^4}\frac{\langle p^2\rangle}{v_R}\frac{(Y_\Delta)_{ee}}{2\rho}\mathcal{M}^{0\nu\beta\beta}(0), \tag{25}$$

where we have used $M_{\Delta_R}^2 \approx 2\rho v_R^2$ and $\mathcal{L} \supset \rho\text{Tr}(\Delta_R\Delta_R^\dagger \Delta_R\Delta_R^\dagger) + Y_\Delta \bar{L}_R^c \Delta_R L_R$.[5] The only possibility of having a phenomenologically relevant contribution is to saturate the bounds on v_R ($M_{W_L}^4/M_{W_R}^4 < 10^{-6}$ [40] and $M_{W_R} \approx v_R g/\sqrt{2}$) having at the same time $Y_\Delta \gg \rho$, which is not very feasible since a small value of ρ would render Δ_R too light, contradicting the experimental bound, $m_{\Delta_R} >$ 320 GeV [59]. We will thus neglect this contribution.

We have shown that only the contributions coming from the light neutrino exchange can have a significant impact on the $0\nu\beta\beta$ decay rate, with the exception of the channel mediated by two W_R gauge bosons in which the heavy neutrino exchange dominates. In summary, the phenomenologically relevant contributions to the $0\nu\beta\beta$ decay rate can be recast as

$$A_{\text{total}} \propto \left[c_{LL}\sum_{i=1}^{3}\left[(U_{\text{pmns}})_{ei}\right]^2 m_i + c_{RR}\left(\frac{M_{W_L}^2}{M_{W_R}^2}+\xi\right)^2 \sum_{i=1}^{3} M_i V_{ei}^{*2}\frac{\mathcal{M}^{0\nu\beta\beta}(M_i)}{\mathcal{M}^{0\nu\beta\beta}(0)} - c_{LR}\,\theta_{e1}^{*}\left(\xi + \eta\frac{M_{W_L}^2}{M_{W_R}^2}\right)\langle p\rangle\right]\mathcal{M}^{0\nu\beta\beta}(0) \equiv m_{\beta\beta}\,\mathcal{M}^{0\nu\beta\beta}(0), \tag{26}$$

[4] Notice that the first and second terms in Eq. (20) correspond to the mechanisms usually called η and λ mechanisms, respectively.

[5] We refer readers to Ref. [21] for more details.

where we have made use of Eq. (6) and c_{LL}, c_{LR} and c_{RR} are coefficients which take into account the different chirality of the outgoing electrons. At this point we can make an estimation of the NMEs associated to the heavy neutrino exchange, $\mathcal{M}^{0\nu\beta\beta}(M_i)$, to understand how relevant the remaining NP contributions are. The effective mass becomes,

$$|m_{\beta\beta}|^2 = \left|\left(\frac{v_L}{v_R}M_R - \theta M_R \theta^T\right)_{ee}\right|^2 + \left|\left(\xi + \eta \frac{M_{W_L}^2}{M_{W_R}^2}\right)\langle p\rangle \theta_{e1}\right|^2 + \left|\left(\frac{M_{W_L}^2}{M_{W_R}^2} + \xi\right)^2 \langle p^2\rangle \left[(M_R)^{-1}\right]_{ee}\right|^2 \quad (27)$$

where we have neglected the suppressed interference terms between the different chirality contributions [60].

In the rest of this section, we will first study the bounds that can be extracted from the $0\nu\beta\beta$ decay experiments if one assumes that the three contributions listed in Eq. (27) are completely independent. After that, we will study the region of the parameter space in which a NP signal in the future $0\nu\beta\beta$ decay experiments can be expected when the correlations among the different contributions are not ignored.

In Fig. 2 we show the constraints on v_R (recall that $M_{W_R} = g v_R/\sqrt{2}$) and the mixing between ν_{eL} and the lightest heavy neutrino, $|\theta_{e1}|$, extracted from $0\nu\beta\beta$ decay experiments when only the contribution from the W_L–W_R channel (second term in Eq. (27)) is taken into account. In the left panel the mixing ξ saturates the theoretical bound ($\xi = M_{W_L}^2/M_{W_R}^2$) while in the right panel ξ is neglected. The shaded region is ruled out by the present constraint, $|m_{\beta\beta}| < 0.38$ eV [9], while the region between the red dashed lines corresponds to the sensitivity of the next-to-next generation of experiments, 10^{-2} eV $< |m_{\beta\beta}| < 0.38$ eV.

Figure 3 is analogous to Fig. 2, but this time we show the present bound on v_R and $(Y_\Delta)_{ee}$ when only the contribution from the W_R–W_R channel is considered, i.e., only the third term of Eq. (27) is included in the analysis. The future sensitivity is shown as well.

The caveat for Figs. 2 and 3 is that we switch on the NP contributions one at a time without considering the correlation between them and that of the light neutrinos. This is specially problematic if one tries to find the future sensitivity to the parameters of the model. For example, from Fig. 2 one would conclude that v_R can be probed in the region 50 TeV $\lesssim v_R \lesssim 4 \cdot 10^4$ TeV, while from Fig. 3 the conclusion would be different, probing 50 TeV $\lesssim v_R \lesssim$ 500 TeV. In this context, the following two questions arise. First, is the standard light neutrino contribution significant for those inputs of the parameters? Can those NP contributions really dominate over the standard one? And second, if yes, for what region of the parameter space? The $0\nu\beta\beta$ decay phenomenology is sometimes analyzed taking into account the different contributions one by one, that is, neglecting the rest of the contributions and the correlations induced by the neutrino mass generation mechanism. In this work, we simultaneously include all the relevant contributions in the analysis and emphasize how the correlation plays a vital role in order to answer the previous questions.

In Fig. 4 we show the sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments (10^{-2} eV $< |m_{\beta\beta}| <$ 0.38 eV) to the parameters of the model by including all the relevant contributions and requiring the NP contribution to the $0\nu\beta\beta$ decay rate (second and third term in Eq. (27)) to be

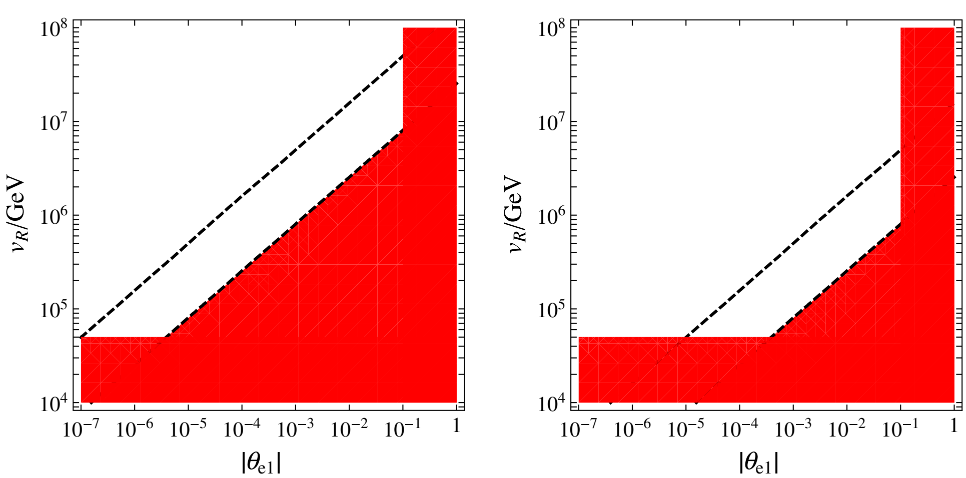

Fig. 2 Heavy regime. The *shaded region* represents the values of v_R and $|\theta_{e1}|$ ruled out by the present experimental bound on the $0\nu\beta\beta$ decay rate mediated by the W_L–W_R channel (neglecting the standard and the W_R–W_R contributions) and the bounds on M_{W_R} [40] and non-unitarity. The future $0\nu\beta\beta$ decay sensitivity, when the standard light neutrino and the W_R–W_R contributions are not included, is given by the region between the *red dashed lines*. The mixing ξ has been fixed to $M_{W_L}^2/M_{W_R}^2$ (zero) in the *left* (*right*) panel

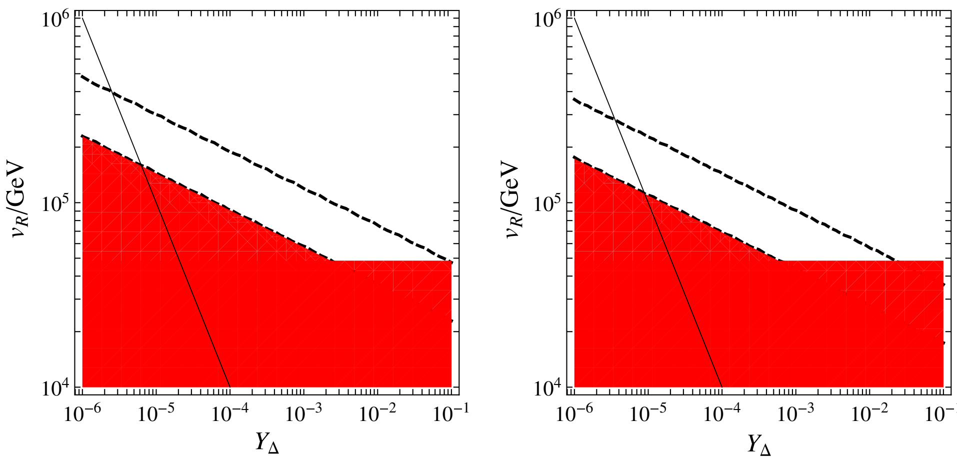

Fig. 3 Heavy regime. The *shaded region* represents the values of v_R and $(Y_\Delta)_{ee}$ ruled out by the present experimental bound on the $0\nu\beta\beta$ decay rate mediated by the W_R–W_R channel (neglecting the standard and the W_L–W_R contributions) and the bounds on M_{W_R} [40]. The future $0\nu\beta\beta$ decay sensitivity, when the standard light neutrino and the W_L–W_R contributions are not included, is given by the region between the *red dashed lines*. The *black line* corresponds to $(M_R)_{ee} = 1$ GeV. The mixing ξ has been fixed to $M^2_{W_L}/M^2_{W_R}$ (zero) in the *left* (*right*) panel

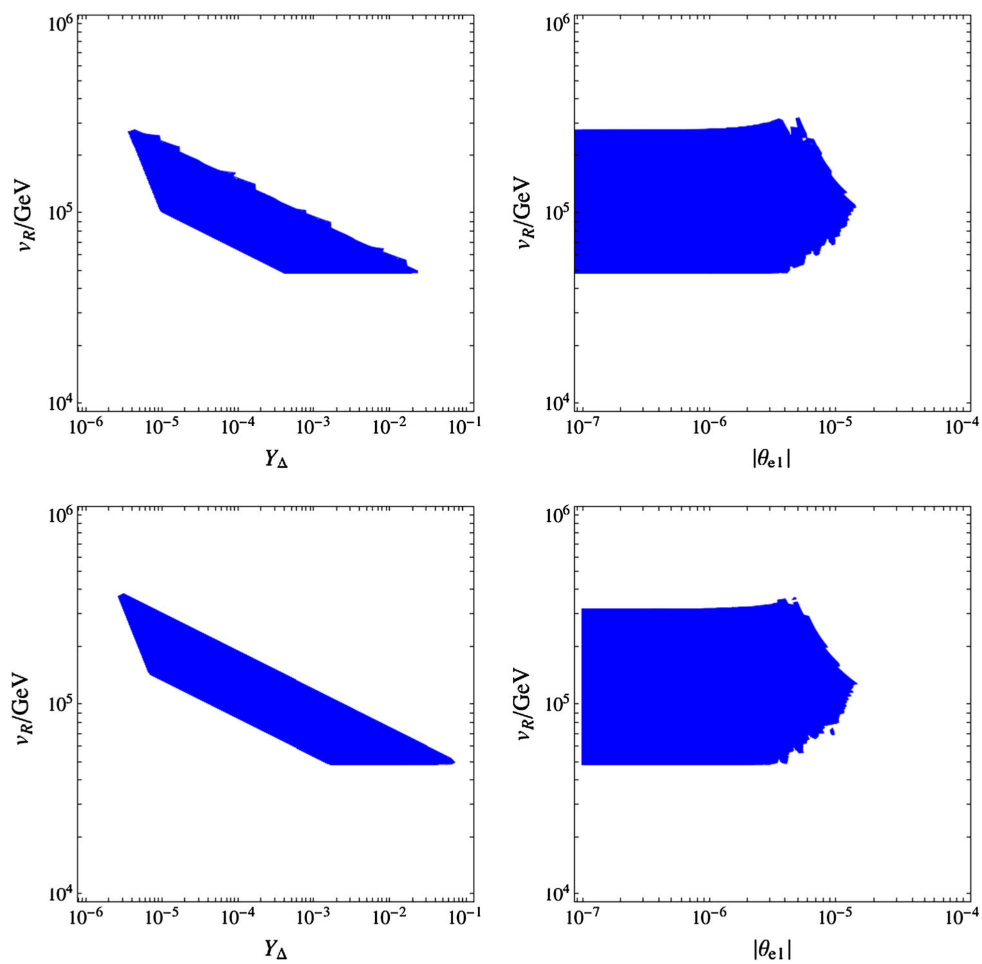

Fig. 4 Heavy regime. The *shaded region* represents the future sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments to v_R and $(Y_\Delta)_{ee}$ (*left panel*) and v_R and $|\theta|$ (*right panel*) when the decay rate is dominated by the NP contribution for $\xi = 0$ (*upper panels*) and $\xi = M^2_{W_L}/M^2_{W_R}$ (*lower panels*). In the calculation all the relevant contributions to the process have been included at the same time in the analysis and the present bound on M_{W_R} [40] and the active–heavy neutrino mixing θ is respected. A hierarchical heavy neutrino spectrum has been considered

at least 10 times larger than the standard contribution (first term of Eq. (27)). The allowed region is projected onto the v_R–$(Y_\Delta)_{ee}$ plane (left panel) and the v_R–$|\theta_{e1}|$ plane (right panel). The mixing has been neglected in the upper panels while in the lower panels is fixed to its maximum value $\xi = M^2_{W_L}/M^2_{W_R}$. The experimental constraints on the W_R

mass [40] and the active–heavy mixing [36,51,52] have also been included. We have assumed that the heavy neutrino spectrum is hierarchical ($M_1 \ll M_2, M_3$). We confirm that a dominant contribution in the left–right symmetric model coming from NP channels is still possible for the window 50 TeV $\lesssim v_R \lesssim$ 300 TeV (50 TeV $\lesssim v_R \lesssim$ 400 TeV) if the trilinear coupling and the mixing are small enough, $3 \cdot 10^{-6} \lesssim (Y_\Delta)_{ee} \lesssim 3 \cdot 10^{-2}$ ($3 \cdot 10^{-6} \lesssim (Y_\Delta)_{ee} \lesssim 8 \cdot 10^{-2}$) and $|\theta_{e1}| \lesssim 2 \cdot 10^{-5}$, respectively, for $\xi = 0$ (maximal mixing $\xi = M_{W_L}^2/M_{W_R}^2$). This corresponds to a range of heavy neutrino masses from GeV to TeV. Comparing the upper and lower panels we can conclude that including the mixing in the analysis has some impact on the results but it is not very significant.

Comparing Fig. 4 with Fig. 2, where only the W_L–W_R contribution is included, we see that the region of the parameter space which can be experimentally probed shrinks when all the contributions are included at once. From Fig. 2, one could conclude that a NP signal from the W_L–W_R channel is possible for very large values of v_R up to $\sim 10^4$ TeV ($M_{W_R} \sim 500$ TeV). However, such a large value of v_R requires a quite large mixing θ since v_R suppresses the W_L–W_R contribution (second term of Eq. (27)) which makes the light neutrino contribution to the $0\nu\beta\beta$ decay rate larger than the present bound. Namely, due to the correlation, such a large values of v_R and θ_{e1} are ruled out and the W_L–W_R channel cannot give a dominant contribution to the process. On the other hand, even for values of v_R close to the present bound, the W_L–W_R contribution is of the same order of the subleading light–active neutrino one, while the W_R–W_R contribution becomes larger than that of the W_L–W_R channel or even above the present experimental bound. Indeed, we have checked numerically that, once the correlations are taken into account, a NP signal can be expected mainly from the W_R–W_R channel as one can anticipate from the fact that the NP signal regions in Fig. 3 and Fig. 4 (left) overlap. Note that the W_L–W_R channel could only dominate the decay rate if a cancelation in the light neutrino contribution takes place, a scenario not explored in this work.

3.2 Light regime

If the RH neutrinos are lighter than the $0\nu\beta\beta$ decay scale, $\mathcal{O}(100)$ MeV, the picture significantly changes with respect to the heavy scenario studied above. Equations (16), (19), and (21) remain valid but the NMEs associated with the "heavy" neutrino exchange are not suppressed compared to the light neutrino mediated ones. In fact, for $M_i < 1$ MeV, we have in a very good approximation $\mathcal{M}^{0\nu\beta\beta}(0) = \mathcal{M}^{0\nu\beta\beta}(M_i)$. This yields a cancelation within the second term of Eqs. (16) and the corresponding amplitude is then given by

$$A_{LL} \propto (M_L)_{ee}\, \mathcal{M}^{0\nu\beta\beta}(0) = \frac{v_L}{v_R}(M_R)_{ee}\, \mathcal{M}^{0\nu\beta\beta}(0), \tag{28}$$

while Eq. (19) becomes

$$A_{RR} \propto \left(\frac{M_{W_L}^2}{M_{W_R}^2} + \xi\right)^2 (M_R^*)_{ee}\, \mathcal{M}^{0\nu\beta\beta}(0), \tag{29}$$

where we have used Eq. (9). The W_L–W_R contribution, given by Eq. (21), vanishes due to the unitarity of the 6×6 neutrino mixing matrix U. Therefore, from Eqs. (28)–(29), the expression for the effective mass $m_{\beta\beta}$ when the RH neutrinos are lighter than 1 MeV becomes

$$\begin{aligned} |m_{\beta\beta}|^2 &= \left|\frac{v_L}{v_R}\,(M_R)_{ee}\right|^2 + \left|(M_R^*)_{ee}\left(\frac{M_{W_L}^2}{M_{W_R}^2} + \xi\right)^2\right|^2 \\ &= \left|v_L\,(Y_\Delta)_{ee}\right|^2 + \left|v_R\,(Y_\Delta^*)_{ee}\left(\frac{M_{W_L}^2}{M_{W_R}^2} + \xi\right)^2\right|^2. \end{aligned} \tag{30}$$

We conclude that in this regime the $0\nu\beta\beta$ decay can be completely attributed to the W_R–W_R channel contribution, if $v_L/v_R \ll \left(M_{W_L}/M_{W_R}\right)^4$.[6]

If $v_L/v_R \gg \left(M_{W_L}/M_{W_R}\right)^4$, the W_L–W_L channel dominates the decay rate. This does not mean that the standard contribution (that mediated by the light–active neutrinos) always dominates since the RH neutrino exchange can also contribute in this channel. Indeed,

$$\begin{aligned} A_{LL} &\propto \left(U_{\text{pmns}}\, m\, U_{\text{pmns}}^T + \theta M_R \theta^T\right)_{ee}\, \mathcal{M}^{0\nu\beta\beta}(0) \\ &= \frac{v_L}{v_R}(M_R)_{ee}\mathcal{M}^{0\nu\beta\beta}(0), \end{aligned} \tag{31}$$

Notice that, contrary to the type-I seesaw limit ($v_L \to 0$) [24], in this regime A_{LL} does not vanish and the RH neutrinos (second term in the equation above) can contribute to the process. Nevertheless, in this work we will focus on the W_R–W_R and W_L–W_R channels. A dominant NP contribution from W_L–W_L channel mediated by the RH neutrinos will be investigated elsewhere in more detail.

The future sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments (10^{-2} eV $< |m_{\beta\beta}| <$ 0.54 eV) to the parameters of the model when the W_R–W_R contribution (second term of Eq. (30)) is at least 10 times larger than that from the W_L–W_L channel (first term of Eq. (30)) is given in Fig. 5. The allowed region of the parameter space is projected this time onto the $v_R - (M_R)_{ee}$ plane (left panel), $v_R - (Y_\Delta)_{ee}$ (central panel) and $v_R - v_L$ (right panel) for $\xi = 0$ (solid line) and $\xi = M_{W_L}^2/M_{W_R}^2$ (dashed line). The bounds on the W_R mass [40] have been also included.

[6] Note that the mass of M_{W_L} mainly comes from the SM Higgs VEV, not from the Δ_L VEV, v_L. As a consequence, v_L could be very small. Furthermore, a small v_L is in better agreement with the $\rho(\equiv M_{W_L}^2/M_{Z_L}^2 \cos^2\theta_W)$ parameter constraints.

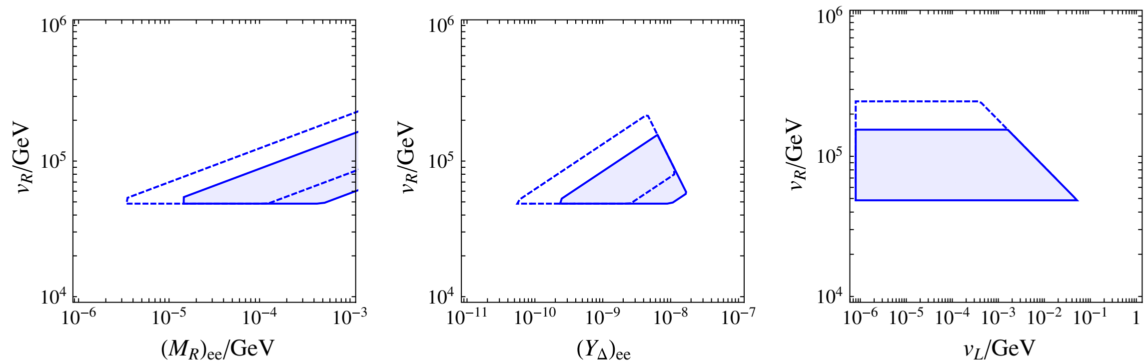

Fig. 5 Light regime. The region inside the *solid* (*dashed*) *lines* represents the future sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments when the decay rate is dominated by the W_R–W_R contribution for $\xi = 0$ ($\xi = M^2_{W_L}/M^2_{W_R}$), projected onto the $v_R - (M_R)_{ee}$ plane (*left panel*), $v_R - (Y_\Delta)_{ee}$ (*central panel*) and $v_R - v_L$ (*right panel*). In the analysis all the relevant contributions have been simultaneously included. The bounds on M_{W_R} [40] have been also included

Figure 5 (left panel) shows that a NP dominated $0\nu\beta\beta$ decay signal can be expected for $(M_R)_{ee} \sim 15$ KeV-1 MeV ($(M_R)_{ee} \sim 3$ KeV-1 MeV) and $M_{W_R} \lesssim 8$ TeV ($M_{W_R} \lesssim 12$ TeV) for $\xi = 0$ ($\xi = M^2_{W_L}/M^2_{W_R}$). This means that in this regime the future sensitivity to M_{W_R} is around a factor 2 weaker than in the heavy regime ($M_{W_R} \lesssim 15 - 20$ TeV depending on the value of ξ). In both regions the sensitivity is driven by the W_R–W_R channel. In order to have a dominant W_R–W_R contribution, the W_L–W_L one should be of course depleted and this can be achieved for small enough values of v_L ($v_L \lesssim 0.07$ GeV) as expected from Eq. (30) and shown in Fig. 5 (right panel). One may ask if it is really feasible or natural to have RH neutrinos lighter than 1 MeV while v_R is above the TeV. Indeed, this is perfectly possible but requires an uncomfortably small value of the trilinear Yukawa coupling Y_Δ since $M_i \sim Y_\Delta v_R$, as shown in Fig. 5 (left). However, the smallness of Y_Δ could be achieved adding an extra mildly broken global symmetry to the model, as has been done in the popular inverse seesaw models with the lepton number. In such a case these small values of Y_Δ could be considered technically natural since $Y_\Delta = 0$ would restore the global symmetry. In any case, it should be remarked that a NP signal can only occur for $10^{-10} \lesssim (Y_\Delta)_{ee} \lesssim 10^{-8}$. Finally, comparing the dashed and solid contours we can conclude that the impact of the mixing ξ is not very significant in this region of the parameter space.

3.3 Mixed scenario

There is an alternative scenario that has not been studied in the previous sections and consists of the existence of RH neutrinos in both regimes below and above the $0\nu\beta\beta$ decay scale. In this section we will focus on the particular case in which one of the RH neutrinos is lighter than 1 MeV and the other two are heavier than 1 GeV, i.e., $M_1 < 1$ MeV and $M_2, M_3 > 1$ GeV, but the phenomenology remains similar if two RH neutrinos are lighter than 1 MeV.

As occurring in the previous section, Eqs. (16), (19), and (21) are also correct in this regime, but only the NMEs associated with the N_2 and N_3 exchange are suppressed compared to the light neutrino mediated ones. The NMEs associated with N_1 satisfy $\mathcal{M}^{0\nu\beta\beta}(0) = \mathcal{M}^{0\nu\beta\beta}(M_1)$. As a consequence, in this regime Eqs. (16) and (21) read

$$A_{LL} \propto \left[\frac{v_L}{v_R}\left(VMV^T\right)_{ee} - \sum_{i=2}^{3} M_i\,(\theta V)^2_{ei}\right]\mathcal{M}^{0\nu\beta\beta}(0), \tag{32}$$

$$A_{LR} \propto -\sum_{i=2}^{3}(\theta^* V^*)_{ei}\,V_{ei}\left(\xi + \eta\frac{M^2_{W_L}}{M^2_{W_R}}\right)\langle p\rangle\mathcal{M}^{0\nu\beta\beta}(0), \tag{33}$$

and Eq. (19) becomes

$$A_{RR} \propto \left(V^{*2}_{e1}M_1 - \sum_{i=2}^{3} V^{*2}_{ei}\frac{\langle p\rangle^2}{M_i}\right) \times \left(\frac{M^2_{W_L}}{M^2_{W_R}} + \xi\right)^2 \mathcal{M}^{0\nu\beta\beta}(0), \tag{34}$$

where again we have used Eq. (9) and the fact that $\mathcal{M}^{0\nu\beta\beta}(M_i)/\mathcal{M}^{0\nu\beta\beta}(0) \ll 1$ for $i = 2, 3$. Therefore, in this scenario the effective mass $m_{\beta\beta}$ is given by

$$|m_{\beta\beta}|^2 = \left|\frac{v_L}{v_R}\left(VMV^T\right)_{ee} - \sum_{i=2}^{3} M_i\,(\theta V)^2_{ei}\right|^2 + \left|\sum_{i=2}^{3}(\theta V)^*_{ei}\,V_{ei}\left(\xi + \eta\frac{M^2_{W_L}}{M^2_{W_R}}\right)\langle p\rangle\right|^2 + \left|\left(V^{*2}_{e1}M_1 - \sum_{i=2}^{3} V^{*2}_{ei}\frac{\langle p\rangle^2}{M_i}\right)\left(\frac{M^2_{W_L}}{M^2_{W_R}} + \xi\right)^2\right|^2. \tag{35}$$

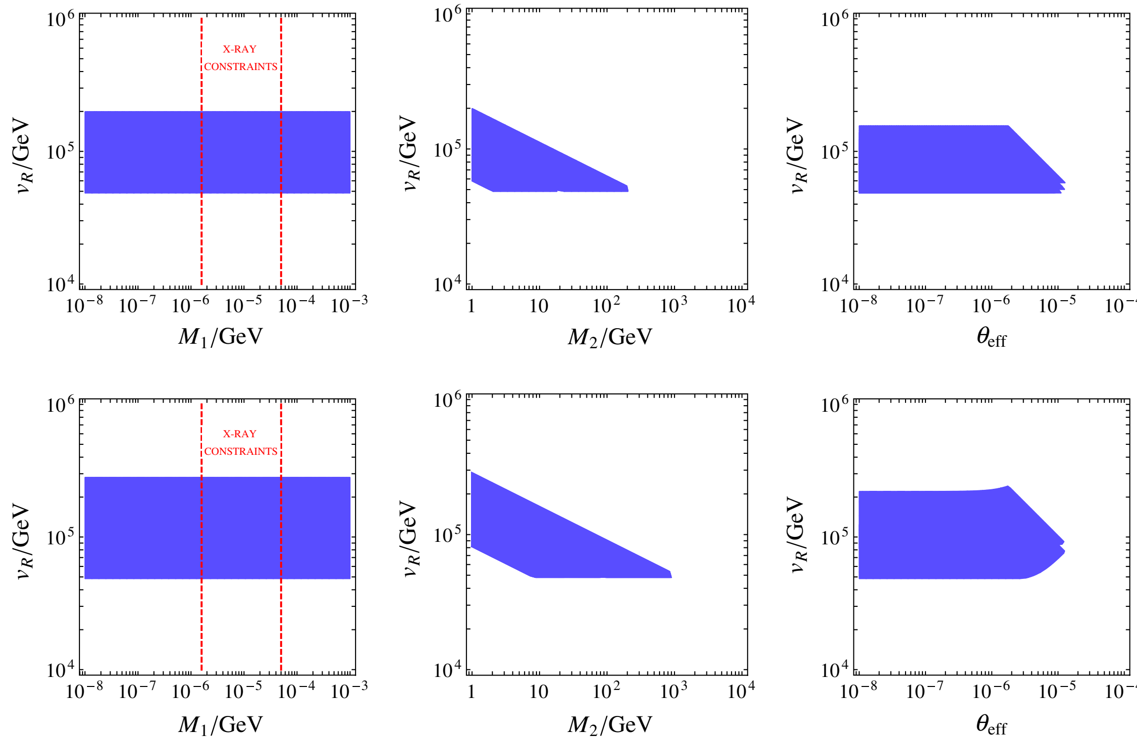

Fig. 6 Mixed scenario. The *shaded region* represents the future sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments projected onto the $v_R - M_1$ plane (*left panel*), $v_R - M_2$ (*central panel*) and $v_R - \theta_{\rm eff}$ (*right panel*) when the decay rate is dominated by the W_L–W_R and W_R–W_R contributions with $\xi = 0$ (*upper panels*) and $\xi = M^2_{W_L}/M^2_{W_R}$ (*lower panels*). In the analysis all the relevant contributions have been simultaneously included. The bounds on M_{W_R} [40], the active–"heavy" mixing [52] and the X-ray constraints [61] have been also included. The PMNS angles and oscillation mass-squared differences have been fixed to the central values given in Ref. [62] and $m_1 = 10^{-2}$ eV, while the CP-phases of $U_{\rm pmns}$ and V have been set to zero

Contrary to the light regime, in this scenario the W_L–W_R contribution may be significant. Notice that if $V = I$, the W_L–W_R contribution cancels out, which means that the RH neutrino mixing V is very relevant in this region.

In order to be consistent with the rest of this work, associated with the absence of cancelation in the light–active neutrino contribution, we will focus in this section on the limit $\frac{v_L}{v_R} \to 0$, neglecting the first term of the W_L–W_L contribution. As expected, we have checked that if that term is switched on in the analysis, cancelations between the two terms of the W_L–W_L contribution can perfectly occur for some part of the parameter space leading to better future sensitivities to the W_L–W_R and W_R–W_R contributions. On the other hand, we will not study the possibility of having a NP signal from the W_L–W_L channel mediated by the RH neutrinos since it would also require some level of fine tuning as demonstrated in Ref. [24] for the type-I seesaw case.

The left and central panel of Fig. 6 show the sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments (10^{-2} eV $< |m_{\beta\beta}| <$ 0.38 eV) to the parameters of the model by including all the relevant contributions and requiring the W_L–W_R and W_R–W_R contribution to the $0\nu\beta\beta$ decay rate (second and third term in Eq. (35)) to be at least 10 times larger than the W_L–W_L contribution (first term of Eq. (35)) and $\xi = M^2_{W_L}/M^2_{W_R}$. The allowed region is projected onto the v_R–M_1 plane (left panel) and v_R–M_2 plane (central panel). We have assumed that $U_{\rm pmns}$ and V are real and fix the light neutrino mass scale to $m_1 = 10^{-2}$ eV. As in the previous plots, the bounds on the W_R mass [40] and the active–"heavy" mixing have been included [36,51,52]. The X-ray constraints [61], which apply if N_1 is the DM, are also shown in Fig. 5: the region between the red dashed lines is ruled out. We have used Eq. (15) in order to be consistent with the light neutrino mass and mixing pattern as mentioned before. However, in order to illustrate better the impact of the mixing θ, we also show in Fig. 6 (right panel) the sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments to v_R and an effective mixing $\theta_{\rm eff}$ defined in the following way. We have assumed in Eq. (35) that $\sum_{i=2}^{3} (\theta V)^*_{ei} V_{ei} \approx \theta_{\rm eff}$ and $\sum_{i=2}^{3} (\theta V)^2_{ei} M_i \approx \theta^2_{\rm eff} M$, which is not true in general but a natural assumption if no particular cancelations are involved. The results for $\xi = 0$ are shown in the upper panels while in the lower panels the W_L–W_R mixing is maximal ($\xi = M^2_{W_L}/M^2_{W_R}$). The mixing plays a role only if it is close to the upper bound and even in that case the future sensitivity is similar to the $\xi = 0$ limit, as can be observed in Fig. 6.

Comparing Figs. 5 and 6, it is clear that the sensitivity to v_R in this case and the light regime is very similar. The main difference comes from the role of the W_L–W_R channel, which in the light regime is completely irrelevant but can be significant in the mixed scenario for values of the W_L–W_R mixing close to the theoretical bound $\xi = M_{W_L}^2/M_{W_R}^2$. In fact, saturating the bound, the W_L–W_R contribution can be the dominant one in the region $\theta_{\rm eff} \sim 10^{-7}$–$10^{-5}$ but for smaller values of $\theta_{\rm eff}$ it is negligible. This is easy to understand since the W_L–W_R contribution is proportional to θ while that of the W_R–W_R channel do not depended on θ. A signal due to the W_L–W_R channel is possible for such a small values of θ thanks to the enhancement from $\langle p \rangle$ due to chirality argument and only if ξ is close to its upper bound. On the other hand, the reason why Fig. 6 shows that the v_R sensitivity is independent of M_1 is because for a large part of the parameter space the W_R–W_R channel dominates, where the N_2 and N_3 contribution is important. In summary, a NP $0\nu\beta\beta$ decay signal can be expected for a quite small light–sterile neutrino mixing ($\theta_{\rm eff} \lesssim 10^{-5}$) if one of the RH neutrinos is lighter than 1 MeV while the rest are above the $0\nu\beta\beta$ decay scale. This is interesting regarding the possibility of accommodating the DM in the left–right symmetric models as we will see in the next section.

4 Complementary constraints

In this section we will study the impact of the one-loop corrections on the light neutrino masses and whether the part of the parameter space which can be probed in future $0\nu\beta\beta$ decay experiments, as described above, is accessible by other experiments.

4.1 One-loop corrections

Since, in the scenarios studied here, the light neutrino contribution to the $0\nu\beta\beta$ decay rate is suppressed with respect to the NP ones, one may expect that this significant NP lepton number violation contribution to the $0\nu\beta\beta$ decay rate could induce non-negligible one-loop corrections to the light neutrino masses. Of course, if the one-loop corrections are larger or similar to the tree-level contribution, they should be included in the analysis, which would modify our previous conclusions. The leading one-loop correction to the light neutrino masses is given by [63,64]

$$(\delta M_L)_{\alpha\beta} = \frac{1}{(4\pi v)^2}\left(\tilde{m}_D^T\right)_{\alpha i}\tilde{M}_i\left\{\frac{3\ln\left(\tilde{M}_i^2/M_Z^2\right)}{\tilde{M}_i^2/M_Z^2-1} + \frac{\ln\left(\tilde{M}_i^2/M_H^2\right)}{\tilde{M}_i^2/M_H^2-1}\right\}(\tilde{m}_D)_{i\beta}\,, \tag{36}$$

where $\tilde{m}_D$ and $\tilde{M} = \mathrm{diag}\,(M_1, M_2, M_3)$ are the Dirac and Majorana sub-matrices, respectively, written in the basis in which the Majorana sub-matrix is diagonal, M_Z is the mass of the Z boson and M_H the Higgs boson mass. Notice that the self energy diagrams with $W_{L,R}$ bosons in the loop do not give any correction to the light neutrino masses since it is proportional to the external momentum [64]. The contribution would have been sensitive to ξ and M_{W_R}, had $W_{L,R}$ contributed to the light neutrino corrections. Assuming that there is no fine tuning and the Yukawa couplings are of the same order, we can roughly estimate the size of the one-loop corrections given by Eq. (36) as

$$\begin{aligned}\delta M_L/m_\nu &\sim 3\left(\frac{M_z}{4\pi v}\right)^2\ln\left(M_i^2/M_Z^2\right) + \left(\frac{M_H}{4\pi v}\right)^2\ln\left(M_i^2/M_H^2\right), \quad \text{for } M_i \gg M_Z, M_H,\\ \delta M_L/m_\nu &\sim \left(\frac{M_i}{4\pi v}\right)^2\left(3\ln\left(M_i^2/M_Z^2\right)+\ln\left(M_i^2/M_H^2\right)\right),\\ &\text{for } M_i \ll M_Z.\end{aligned} \tag{37}$$

Using the estimation given by the first equation above, we can conclude that, if $M_i \gg M_Z, M_H$, the one-loop corrections to the light neutrino masses are under control for the range of values that can be probed in future $0\nu\beta\beta$ decay experiments. In fact, for $M_i \sim 1$ TeV we have $\delta M_L/m_\nu \sim 10^{-2}$, and $\delta M_L/m_\nu$ gets smaller for smaller values of M_i; for example, $\delta M_L/m_\nu \ll 10^{-4}$ for $M_i \ll M_Z$. This can be understood as follows. From Eq. (36), one can infer that the tree-level contribution is bigger than the loop induced ones because it has a similar structure to $m_D M_i^{-1} m_D^T$ but without the loop suppression, $1/(16\pi^2)$. This is correct unless some cancelation is at work for the tree-level contribution, which is not the case studied here. Notice that, in this sense, the assumptions made in order to obtain Eq. (37) are quite reasonable.

Therefore, we can conclude that the one-loop corrections to the light neutrino masses are negligible and not relevant in our analysis. The lepton number violation source of the dominant NP contributions studied in the previous sections is the Majorana mass term generated dynamically for the RH neutrinos. Indeed, this source of lepton number violation is related to the light neutrino masses through the seesaw mechanism, and this correlation has been taken into account in the previous analysis. The dominant NP contribution to the $0\nu\beta\beta$ decay coming from the W_R–W_R channel (or the W_L–W_R channel in the mixed scenario) requires the suppression of the standard (and long range) light neutrino one. Since we are not facing the possibility of having any cancelation in the light neutrino contribution, in order to achieve this suppression the Yukawa couplings and v_L/v_R should be small. The W_R–W_R contribution can be dominant because the RH mixing is not constrained in contrast with the active–heavy mixing θ, which is necessarily small as the Yukawa couplings. The W_L–W_R channel can dominate in the mixed scenario (only for large ξ) due to the enhancement coming from the NME and the linear dependence on the active–heavy mixing θ.

4.2 Other experimental bounds

The charged LFV experiments are also sensitive to the parameters of the model that can be probed in $0\nu\beta\beta$ decay experiments. Among them, $\mu \to e\gamma$, $\mu \to 3e$ and $\mu \to e$ conversion give the stronger bounds. First of all, the small active–heavy mixing required here in order to have a significant NP contribution to $0\nu\beta\beta$ decay ($\theta \lesssim 10^{-5}$), renders the type-I seesaw like contribution to this processes completely negligible since the strongest present bound coming from $\mu \to e\gamma$ gives the constraint $|(\theta^\dagger\theta)_{e\mu}| \lesssim 10^{-5}$. A complete calculation of the charged LFV branching ratios in the MLRSM can be found in [29]. The most relevant constraint in the context of this work comes from $\mu \to e\gamma$, whose branching ratio, to zeroth order on θ and ξ, is given by

$$Br_{\mu\to e\gamma} \approx 2.6\times 10^{-10}\left(\frac{\text{TeV}}{M_{W_R}}\right)^4\left(\frac{|\left(M_R M_R^*\right)_{\mu e}|}{M_{W_R}^2}\right)^2, \tag{38}$$

for $M_{\Delta_{L,R}} \gg M_{W_R}$. Applying the present experimental constraint [65] to Eq. (38), the bound on M_R reads

$$M_R \lesssim \left(\frac{M_{W_R}^2}{4.6\ \text{TeV}}\right). \tag{39}$$

Saturating the lower bound on M_{W_R}, one obtains $M_R \lesssim$ 1 TeV. This is the largest M_R that can be probed with $0\nu\beta\beta$ decay experiments as can be seen in Fig. 4 (it corresponds to the bottom right corner of the shaded regions in the left panels). This means that the future $\mu \to e\gamma$ experiments can be sensitive at least to that corner of the parameter space, which can also give a signal in $0\nu\beta\beta$ decay experiments. One should, however, keep in mind that the flavor structure of M_R plays an important role, being $\mu \to e\gamma$ experiments indeed sensitive to $\left(M_R M_R^*\right)_{\mu e}$ and $0\nu\beta\beta$ decay mainly to $(M_R)_{ee}$ if the heavy neutrino spectrum is hierarchical. Therefore, a NP signal in $0\nu\beta\beta$ decay experiments does not necessarily imply also a signal in future $\mu \to e\gamma$ experiments. On the other hand, the bound in Eq. (39) has been extracted assuming that $M_{\Delta_{L,R}} \gg M_{W_R}$, but smaller masses of the triplets can clearly enhance the branching ratio [29]. The same applies for the $\mu \to 3e$ and $\mu \to e$ conversion case since their branching ratios are inversely proportional to the triple masses. Therefore, we cannot extract a bound like Eq. (39) from $\mu \to 3e$ and $\mu \to e$ conversion, since for large triplet masses the branching ratios are very suppressed.

So far, in our LFV analysis we have neglected the ξ contribution. If one switches on the left–right mixing ξ, the following constraint from $\mu \to e\gamma$ can be extracted [33]:

$$|\,(m_D)_{\mu e}\,|\xi \lesssim 2\ \text{KeV}, \tag{40}$$

which can be roughly translated into $\theta\xi \lesssim \frac{2\ \text{KeV}}{M_R}$, which is basically compatible with most of the parameter space that can give a NP signal in $0\nu\beta\beta$ decay, since $\xi < 10^{-3}$ and $\theta < 10^{-5}$. Only if M_R is close to TeV and ξ saturates the present bound ($\xi \le M_{W_L}/M_{W_R} < 10^{-3}$), then could the mixing ξ have an impact on $\mu \to e\gamma$. Basically, we could probe the same part of the parameter space as commented on above but for values of ξ close to its present bound, again with the important warning that the flavor structure plays an essential role here.

Finally, in the MLRSM the electric dipole moment (EDM) can be considerably enhanced with respect to the SM result (up to 10 orders of magnitude). This is because the SM contribution to the EDM appears at four loops while the left–right symmetric model can provide a huge enhancement due to the left–right mixing ξ [66]. In fact, the EDM experiments can be sensitive in the future to part of the parameter space studied here [31], mainly through the imaginary part of $\left[(m_D)_{ee}\,\xi\right]$.

5 Dark matter

In this section we study the possibility of having a successful DM candidate in the context of the MLRSM when the $0\nu\beta\beta$ decay rate is dominated by NP contributions. The first question which arises from the results of the previous section is whether N_1 can be DM in the light regime, namely with mass $M_1 \lesssim \mathcal{O}(\text{MeV})$. This reminds us of the Dodelson–Widrow (DW) scenario [67], where a KeV neutrino is produced via neutrino oscillations and can be a viable DM candidate.[7] In the left–right symmetric models, however, a RH KeV neutrino N_1 would be thermally produced via the W_R or Z_R exchange and decouples from the thermal bath at the freeze-out temperature T_f,

$$T_f \sim 400\ \text{MeV}\left(\frac{g_*(T_f)}{70}\right)^{1/6}\left(\frac{M_{W_R}}{5\ \text{TeV}}\right)^{4/3}, \tag{41}$$

where $g_*(T_f)$ is the number of relativistic degrees of freedom at freeze-out. The rule of thumb to estimate T_f is to set the interaction rate equal to the expansion rate of the Universe. Given that we are interested on the region of the parameter space in which the NP dominates the $0\nu\beta\beta$ decay rate, the W_R mass should be in the range $M_{W_R} \sim$ 1–15 TeV (see Fig. 6).

[7] A recent study for KeV-neutrino DM on the $0\nu\beta\beta$ decay in the context of the type-I seesaw can be found in Ref. [68], where the KeV-neutrino contribution to the $0\nu\beta\beta$ rate is subleading due to the X-ray bound.

Therefore, N_1 is highly relativistic ($M_1 \lesssim$ MeV $\ll T_f$) at freeze-out and the resulting relic density is [69]

$$\Omega_{N_1} \simeq 3.3 \left(\frac{M_1}{1\ \text{KeV}}\right)\left(\frac{70}{g_*(T_f)}\right), \tag{42}$$

which, for $M_1 \sim$ KeV, would be much larger than the observed DM relic density $\Omega_{\text{DM}} = 0.265$ [70]. This constraint is much severer than the X-ray constraints shown in Fig. 5 and the Big Bang Nucleosynthesis (BBN) bound, which is indeed still compatible at $\sim 2\sigma$ with the existence of one extra relativistic species [71,72]. A possible way out has recently been proposed and studied in detail in Refs. [69,73]. Basically, the idea is to dilute the number density of N_1 by the injection of entropy into the thermal bath after N_1 freezes out. To be more specific, the out-of-equilibrium decay of N_2 and/or N_3, of mass around GeV, into SM particles can increase the entropy of the Universe, leading to faster Universe expansion and in turn a smaller N_1 density. The set of constraints that should be satisfied if N_1 as DM was once in thermal equilibrium has been summarized in Ref. [73]. In particular, the authors claim that the required entropy injection can be achieved if $M_{W_R} \gtrsim$ 10–16 TeV, while the Lyman-α constraints require $M_1 \gtrsim 1$ KeV. It turns out that these bounds and the rest of the constraints listed in Ref. [73] are compatible with a future NP signal in $0\nu\beta\beta$ decay experiments described in Sect. 3.3. Notice that the RH neutrino spectrum required to have DM ($M_1 \sim$ KeV and $M_2, M_3 \sim$ 1–10 GeV) belongs to the mixed scenario where a NP signal in $0\nu\beta\beta$ decay experiments is possible. It should be remarked that the constraint $M_{W_R} \gtrsim$ 10–16 TeV is in obvious tension with a future NP signal in the $0\nu\beta\beta$ decay. However, as we have mentioned in the previous section, if v_L is switched on in the above analysis a cancelation between the two terms in the light neutrino contribution can take place such that a NP signal is possible for $M_{W_R} \gtrsim 16$ TeV.

On the other hand, in Ref. [69] an alternative scenario able to relax the bound on M_{W_R} from Ref. [73] was carefully analyzed. In this scenario the desired dilution of the number density of N_1 is achieved for $M_1 \simeq 0.5$ KeV, $M_2 \sim 140$ MeV, and $M_3 \sim 245$ MeV, with $M_{W_R} \sim 5$ TeV and the help of a particular right-handed flavor structure such that N_2's coupling constant to SM leptons is stronger than that of N_1. We refer the readers to Ref. [69] for the details of the analysis. In any case, as was already pointed out in Ref. [69], the contribution to the $0\nu\beta\beta$ decay rate from the W_R–W_R channel associated with such spectrum can be testable in the future $0\nu\beta\beta$ decay experiments as we have confirmed in Sect. 3.3. However, we would like to remark that the W_L–W_R contribution can also be very relevant in this case, as was explained in the previous section.

Finally, in the left–right symmetric models, in principle the neutral component of the right-handed triplet Δ_R^0, which is a singlet under the SM gauge group, could also be a DM candidate. However, it decays at one-loop into two photons via W_R exchange [69], i.e.,

$$\Gamma_{\Delta_R^0\to\gamma\gamma} \sim 10^{-52}\ \text{GeV}\left(\frac{m_\Delta}{\text{KeV}}\right)^3\left(\frac{10^{13}\ \text{GeV}}{M_{W_R}}\right)^2. \tag{43}$$

The X-ray constraints on KeV DM resulting from observations on galaxies and clusters of galaxies [74] requires $\tau_{\Delta_R^0\to\gamma\gamma} = 1/\Gamma_{\Delta_R^0\to\gamma\gamma} \gtrsim 10^{28}$ s or $\Gamma_{\Delta_R^0\to\gamma\gamma} \lesssim 10^{-52}$ GeV. Therefore, a KeV Δ_R^0 would imply a too heavy W_R such that the contribution of the NP channels involving W_R to the $0\nu\beta\beta$ decay would be completely negligible. Nevertheless, the X-ray constraints apply only to DM with masses around 1–20 KeV. In principle this leaves another window of Δ_R^0 mass which can be studied. However, other constraints make this possibility quite unfeasible. First, the mass of DM is constrained to be larger than KeV [75] because of the Lyman-α observations. Second, for $M_{\Delta_R^0} \gtrsim 20$ KeV, $\tau_{\Delta_R^0}$ still has to be longer than the age of the Universe, around 10^{18} s, which results again in a very heavy W_R that renders any NP contribution to $0\nu\beta\beta$ decay far beyond the future experimental sensitivity.

In summary, in spite of the existence of various constraints, the left–right symmetric models can accommodate a KeV RH neutrino as a successful DM candidate which can lead to a NP signal in the next-to-next generation of $0\nu\beta\beta$ decay experiments.

6 Conclusions

We have studied the $0\nu\beta\beta$ decay phenomenology in the MLRSM. In particular, we have analyzed under which conditions a $0\nu\beta\beta$ decay signal can come mainly from NP contributions associated with this model. Special attention has been paid to the correlation among the different NP contributions and the standard light neutrino one. This correlation emerges from the neutrino mass generation mechanism and should always be considered in the analysis. The scenario in which an accidental cancelation in the W_L–W_L contribution takes place has not been explored. The role of the W_L–W_R mixing ξ and the possibility of having a NP dominated $0\nu\beta\beta$ decay signal compatible with DM is also investigated.

We have distinguished three different regions of the parameter space based on the mass of the RH neutrinos: (i) all the masses heavier than GeV, denoted by heavy regime; (ii) masses lighter than MeV, dubbed light regime; (iii) the lightest mass below the MeV and the rest above GeV, called mixed scenario. Notice that (i) has been extensively studied in the literature, but (ii) and (iii) have not been analyzed before in detail for the left–right symmetric models (at least the $0\nu\beta\beta$ decay phenomenology).

In the heavy region, we have found that the dominant NP contribution emerges mainly from the W_R–W_R channel mediated by the heavy neutrinos. To be more precise, it has been shown that this dominant NP contribution could be measured in the next-to-next generation of $0\nu\beta\beta$ decay experiments for the window $M_{W_R} \sim 1 - 15$ TeV ($M_{W_R} \sim 1 - 20$ TeV), if the active–heavy mixing is smaller than $\sim 10^{-5}$ and the right-handed triplet "Yukawa" coupling satisfies $3 \cdot 10^{-6} \lesssim (Y_\Delta)_{ee} \lesssim 3 \cdot 10^{-2}$ ($3 \cdot 10^{-6} \lesssim (Y_\Delta)_{ee} \lesssim 8 \cdot 10^{-2}$), which corresponds to a range of heavy neutrino masses from GeV to TeV for $\xi = 0$ ($\xi = M^2_{W_L}/M^2_{W_R}$). We have also shown that neglecting the present correlation between the various contributions, and that with the light neutrinos in particular, can lead to incorrect results. For instance, it is found that the region of the parameter space which can be experimentally probed when only the W_L–W_R contribution is included in the analysis shrinks considerably if all the contributions are included at once and their correlations are not neglected.

The results for the light region turn out to be similar to those of the heavy region. We have found that a future $0\nu\beta\beta$ decay NP signal can come only from the W_R–W_R channel since the W_L–W_R contribution cancels out in this regime. In particular, we have shown in which part of the parameter space this is possible and we found a similar but weaker sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments: $M_{W_R} \lesssim 8$ TeV ($M_{W_R} \lesssim 12$ TeV) for $\xi = 0$ ($\xi = M^2_{W_L}/M^2_{W_R}$). A NP signal can be expected for $(M_R)_{ee} \sim 15$ KeV−MeV ($(M_R)_{ee} \sim$ KeV−MeV) for $\xi = 0$ ($\xi = M^2_{W_L}/M^2_{W_R}$), while the triplet Yukawa coupling should be inside the region $10^{-10} \lesssim (Y_\Delta)_{ee} \lesssim 10^{-8}$. This uncomfortably small value of Y_Δ is required in order to have very light RH neutrinos ($M_i < 1$ MeV) since $M_i \sim Y_\Delta v_R$; that seems unnatural but might be achieved with the help of an additional global symmetry. On the other hand, notice that in this regime the W_L–W_L contribution is proportional to v_L/v_R and therefore a small value of v_L ($v_L \lesssim 0.07$ GeV) guarantees a dominant W_R–W_R contribution to the $0\nu\beta\beta$ decay rate. If $v_L/v_R \gg (M_{W_L}/M_{W_R})^4$, which is still experimentally allowed, the RH neutrinos can dominate the process via the W_L–W_L channel, contrary to the type-I seesaw case in which the decay rate is very suppressed if all the RH neutrinos are lighter than the $0\nu\beta\beta$ decay scale.

In the intermediate regime, with RH neutrinos in both regions ($M_1 \lesssim 1$ MeV and $M_2, M_3 \gtrsim 1$ GeV), if the W_L–W_R mixing is close to the theoretical upper bound $\xi = M^2_{W_L}/M^2_{W_R}$, the role of the W_L–W_R channel can be relevant, in contrast with the previous cases. We have found that a NP signal coming from the W_L–W_R channel could take place for $M_{W_R} \sim$ 1–10 TeV and an active–heavy neutrino mixing $\theta \sim 10^{-7}$–10^{-5}. Indeed, a signal from the W_R–W_R channel can be expected in a larger region, even for smaller values of θ since its contribution is independent of the active–heavy neutrino mixing. In this case we have focused our study on the limit $v_L/v_R \to 0$, but we have checked that if v_L is switched on in the analysis the future sensitivity to v_R is much better. However, this can only take place when there is a cancelation between the type-I and type-II seesaw terms in the light neutrino contribution.

In order to study the impact of the W_L–W_R mixing ξ, we have analyzed the following two extreme limits: $\xi = M^2_{W_L}/M^2_{W_R}$ (maximal) and $\xi = 0$ (negligible). We have shown that the inclusion of the W_L–W_R mixing ξ can have some impact on the results but it is not very significant, with the possible exception of the mixed scenario where a large mixing is required in order to have a relevant role of the W_L–W_R channel. In general, due to the enhancement on the W_R–W_R contribution for maximal mixing, the sensitivity to M_{W_R} is about a factor 1.5 larger for $\xi = M^2_{W_L}/M^2_{W_R}$ than for $\xi = 0$ in all the regions under study.

We have also analyzed the role of the complementary bounds coming from charged LFV processes and the impact of the one-loop corrections to the light neutrino masses in the context of this work. It turns out that the light neutrino masses are stable under one-loop corrections since they might be important only if a cancelation takes places in the light neutrino masses, but not in the case studied here where there is a general suppression of the light masses with small v_L/v_R and the Yukawa couplings. Future charged LFV experiments might allow us to probe part of the parameter space that can be responsible for a NP signal in $0\nu\beta\beta$ decay experiments, but only a small region in the heavy regime around $M_R \sim 1$ TeV (the bottom right corner of the shaded regions in the left panels of Fig. 4). In fact, a more complete study, beyond the scope of this work, including the effect of triplet masses close to their lower bounds, which can enhance the branching ratios, would be required in order to clarify the issue. A large left–right mixing ξ can also be probed in future EDM experiments as was shown in Ref. [31].

Finally, the following DM-related question has also been addressed. Can a NP dominated $0\nu\beta\beta$ decay signal be compatible with a successful DM candidate in the left–right symmetric models? We conclude that, regardless of the various strong constraints, it is still possible for the scenario proposed in Ref. [73], where a KeV RH neutrino can be the DM if the scale of the other heavy neutrinos is around 1–10 GeV and $M_{W_R} \gtrsim$ 10–15 TeV. We have shown that the $0\nu\beta\beta$ decay signal can be induced by the RH neutrinos through the W_L–W_R and W_R–W_R channel. Additionally, Ref. [69] opens a window of $M_{W_R} \sim 5$ GeV within the horizon of LHC after the QCD phase transition is carefully included.

Acknowledgments We thank to S. Petcov and Javier Menendez for useful discussions and important remarks. This work was partially supported by the ITN INVISIBLES (Marie Curie Actions, PITN- GA-2011-289442). W.-C.H. would like to thank the hospitality of IFPA at Université de Liège, where part of this work was performed.

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